Constraints on the Existence of Strange Quark Stars

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Creation of strange quark stars through strong interaction deconfinement is studied based on modern estimates of hyperon formation in neutron stars. The hyperon abundance is shown to be large enough so that if strange quark matter (SQM) is the true ground state of matter, the deconfinement density should be at most 2.5–3 times the nuclear saturation density. If so, deconfinement occurs in neutron stars at birth, and *all* neutron stars must be strange quark stars. Alternatively, sould observation indicate that some neutron stars have a baryonic interior, SQM is unlikely to be absolutely stable.

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One of the implications of the hypothesis that strange quark matter (SQM) is the true ground state of matter [1,2] is that some or all of neutron stars are actually strange quark stars. The properties of strange quark stars have been studied in many works [3-6], and were found to be mostly similar to those of "conventional" neutron stars, where matter is in a baryonic phase. Recent studies of strange quark stars include unique cooling scenarios [7] and formation of strange dwarfs [8].

A key issue regarding strange star formation is the conversion (deconfinement) process of baryonic matter to SQM. A direct consequence of the SQM hypothesis [1] is that if a seed of SQM is created in the interior of a neutron star, it will initiate a burning process that converts the entire star within a time scale of minutes [9-11]. It has been recently proposed that the emitted energy in such a conversion, typically estimated at several times 10^{52} ergs, could be a source of γ-ray bursts of cosmological origin [12].

Several mechanisms for creating SQM seeds in neutron stars have been suggested [3,10]. These can be either internal, when the SQM seed is created in the star through hadron deconfinement, or external, when a SQM "nugget" which has been created elsewhere is absorbed in the star (or earlier in its progenitor). The abundance of free SQM nuggets in the universe (for example, from debri of strange stars which have coalesced with a binary counterpart) could be large [13], but is diffuclt to evaluate and initially depends on creation of strange quark stars through internal mechanisms. On the other hand, availability of internal mechanisms depends only on the structure of the neutron star. Furthermore, assuming the SQM hypothesis is correct, then if some internal mechanism of SQM seed formation is available in neutron stars at birth, all existing neutron stars must actually be strange quark stars. Alternatively, neutron stars can serve as a test of the SQM hypothesis: if a corrseponding internal mechanism is indeed available, but some of the observed neutron stars can be ruled out as being strange quark stars, the likely conclusion is that SQM is not the ground state of matter.

The purpose of this Letter is to illuminate the role of hyperon formation as an available mechanism for creating SQM in all neutron stars *at birth*. A SQM seed can form within the neutron star if deconfinement of the hadronic matter is energetically favorable. Such deconfinement must proceed through the strong interaction, since multiple creation of strange quarks through the weak intearaction is suppressed. Hyperons, along perhaps with K^- meson condensation, can initiate an internal mecahnism to create the SQM seed by providing the necassary strange quark content so that deconfinement proceeds through the strong interaction. While K^- condensation is typically found to occur only at considerably higher densities of baryonic matter in equilibrium [14,15], the consensus amoung recent works is that hyperon formation in neutron stars should begin at rather low densities. In fact, it is widely accepted [15-18] that hyperons begin to accumulate at about twice the nuclear

saturation density, ρ_0 ($\rho_0 \approx 0.16$ fm^{-3}). An example of the composition of neutron star matter in equilibrium, i.e., the fraction of each species as a function of the total baryon density, is given in Figure 1 (based on pevious work [16]).

At low densities ($\rho_B \leq 3\rho_0$) the relevant hyperons are the Σ^- and Λ. Their combined abundance builds up a strangeness per baryon fraction, $|S|/A$, that exceeds 0.1 at $\rho_B \approx 2.5\rho_0$, and at $\rho_B \approx 3\rho_0$ is already $\sim$ 0.2. These trends result from employing realistic hyperon$-$nuclear-matter interactions based on hypernucler data, and are weakly dependent on the model and the corresponding estimated equation of state.

Contrary to claims in some previous works [10], hyperon induced deconfinement does not require the baryon strangeness fraction to be equal to that of ground state SQM, $|S|/A \approx 0.7 - 0.8$ [2]. Rather, the condition is only that the phase transition into SQM with a composition identical to that of the baryonic matter, e.g., strong deconfinement, be energetically favorable; the SQM then reaches its ground state by series of weak decays.

It is often assumed that the strong interaction deconfinement requires the entire bulk to deconfine spontaneously into quark matter of equal composition and baryon number density (a discussion of a transition through an intermediate mixed phase follows below). In this case, the bulk baryon density, ρ_B, at which the transition is expected satisifies the condition

$$\frac{\varepsilon}{\rho_B}\left(\{b_i\}, \rho_B\right) = \frac{\varepsilon}{\rho_B}\left(\{q_i(\{b_i\})\}, \rho_B\right) \quad , \tag{1}$$

where ε is the energy density and ε/ρ_B is the corresponding energy per baryon, $\{b_i\}$ denotes the baryon equilibrium composition and $q_i(\{b_i\})$ is the appropriate deconfined composition. If condition (1) is satisfied, a spontaneous strong interaction transition into SQM will occur. Whether or not deconfinement will occur in a neutron star depends on the resulting density of deconfinement. If this density, which is dependent on the equations of state of the two phases, is reached in the star's interior, deconfinement is expected.

The quark matter equation of state may be evalulated within the MIT bag model detailed in [2]. The specifics of the bag model are determined by the combinations of values for the bag constant, B, and the quark interaction coupling constant, α_c. For each value of α_c, B is limited from below so that two-flavor quark matter is less bound than symmetric nuclear matter, since ordinary nuclei do not deconfine strongly. The SQM hypothesis places an upper limit on B, in order for the ground state composition of SQM to be more bound than symmetric nuclear matter. Combinig these two conditions constrains the allowed range of values for the bag constant for any given value of α_c which is consistent with the SQM hypothesis [2]. For example, for $\alpha_c = 0$, B is limited to the range 56 MeV fm$^{-3} \leq B \leq$ 82 MeV fm^{-3}.

In order to estimate the density of deconfinement, energies per baryon were calculated for the baryonic and quark phases using various models of the baryonic equation of state

from [16], and different bag models for the quark equation of state. As an example, Figure 2 compares the energy per baryon of baryonic matter with hyperons calculated with a model similar to $\delta=\gamma=\frac{5}{3}$ of [16], to the energy per baryon of quark matter with identical quark composition, using different combinations of B and α_c. The u and d quarks were assumed to be massless, and the mass of the s quark is set to 150 MeV. The transition density for each combination corresponds to the point where the energy per baryon in the baryonic phase crosses that of the quark phase.

The nonzero fraction of strange quarks lowers the energy per baryon in the deconfined phase, and it is found in this work that *all* models which predict that SQM is the true ground state also predict that spontaneous deconfinement should occur at densities lower than $3\rho_0$. It can be seen that a deconfinement density of $\rho_B \leq 2.5\rho_0$ is found even for combinations of B and α_c which are borderline for making SQM more bound than ordinary nuclear matter, i.e. ($B=82$, $\alpha_c=0$), or ($B=63$, $\alpha_c=0.3$) [2]. Even in the combination ($B=100, \alpha_c=0$), for which the SQM hypothesis is no longer correct, the deconfinement density is still found to be lower than $3\rho_0$.

Quark matter models which do not allow spontaneous deconfinement are still possible, of course, like ($B = 125$, $\alpha_c = 0$) shown in Figure 2. In such models the corresponding binding energy of SQM is significantly larger than ordinary nuclear matter. In this case, some deconfinement might occur, and the resulting state could be coexisiting baryon and quark phases [19,20].

The low values for the deconfinement density are mostly due to the presence of strange quarks. If hyperon formation is ignored, the baryonic matter includes only two flavors of quarks. It is found that nuclear matter in beta equilibrium may deconfine into two-flavor quark matter if quark matter is very bound (B must be close to its lower limit for a given value of α_c). For extreme quark matter models (such as $B=56$ MeV fm^{-3}, $\alpha_c=0$), deconfinement may occur even at densities of $\rho_B \approx \rho_0$, since nuclear matter at beta equilibrium has a higher energy per baryon than symmetric matter. However, other combinations of B and α_c delay the deconfinement of nuclear matter to higher densities, and in some cases, quark matter does not form at any density.

As is expected, Similar results are found for other models of the baryon equation of state, since the equation of state of matter with hyperons is limited to a rather narrow range of values [16]. In any case a valid equation must be stiff enough to support a maximum mass of at least 1.4 $M_\odot$ (the determined mass of the 1913+16 pulsar). While the quark bag model is clearly a simplified description of quark matter physics, it seems that the margin it allows for low density deconfinement into SQM is large enough, so that the qualitative results are unlikely to be dependent on the quark matter model as well.

The baryonic density at which the bulk deconfines may actually be regarded as an upper

limit for the creation of quark matter. This may be demonstrated by considering an alternative scenario, where quark matter drops form in the matter through quasi-equilibrium combinations of coexisting baryonic and quark phases. Once finite size quark phase bubbles appear they act as a seed of SQM, which burns into its equilibrium composition and consumes the surrounding baryons through further weak interactions.

Equilibrium between baryonic and quark phases invovles conservation of two charges (baryon number and electric charge), and so the phases need not have equal compositions nor equal densities. This has been pointed out by Glendenning [19] with respect to two phases in full equilibrium, and is also valid for quasi-equilbrium. The quasi-equilibrium conditions differ from those for two phases in full equilibrium [19,20] since the initial deconfinement is assumed to be determined by the strong interaction alone (again, multiple creation of strange quarks is forbidden). For matter with a given composition and total baryon number density the conditions for two-phase equilibrium are the appropriate Gibbs conditions for chemical and pressure equilibrium:

$$\begin{aligned} \mu_n = 2\mu_d + \mu_u \quad &, \quad \mu_p = 2\mu_u + \mu_d \\ \mu_\Lambda = \mu_d + \mu_u + \mu_s \quad &, \quad \mu_{\Sigma^-} = 2\mu_d + \mu_s \end{aligned} \tag{2}$$

and

$$P_B = P_Q \quad , \tag{3}$$

where μ_i refers to the chemical potential of species i, and P_B and P_Q are the pressure in the baryonic and quark phase, respectively. Since weak interactions are ignored, full equilibrium between the species is not enforced, and each baryon species equilibriates with the quark phase seperately. Furthermore, during deconfinemnt the total number of quarks of each of the three species must remain constant. Deconfinement proceeds once equilibrium allows the quark phase to occupy a nonzero fraction of the volume.

In this work deconfinement through quasi-equilibrium was found to occur at lower baryon densities than bulk deconfinement, mainly because the quark phase can have a differnet density and composition than the baryonic phase. For combinations of B and α_c which enable SQM to be absolutely stable, coexisting baryonic and quark phase are found to appear even at $\rho_B = \rho_0$. This is in agreement with the results of [19,20] regarding full equilibrium. In fact relaxing the condition of beta equilibrium in the quark phase yields even slightly lower densities of deconfinement than for full equilibrium. Appropriately, the bag model constants which prevent low density deconfinement are even further away from the range where SQM is absolutely stable: for example, for $\alpha_c = 0$, only $B \approx 200$ can delay deconfinement to $3\rho_0$. While quasi-equlibrium deconfinement might be suppressed (for example, due to finite-size effects), these results offer supprot to the main conclusion of

the analysis of bulk deconfinement: if the SQM hypothesis is correct, deconfinement should occur at densities below $2.5-3\rho_0$.

The immdediate consequence of the above discussion is that in neutron stars with a central density greater or equal to $2.5-3\rho_0$, the baryons should deconfine into SQM, if such matter is indeed absolutely stable. The SQM will then convert through the weak interaction to its equilibrium composition, and proceed to convert the entire star into a strange quark star.

Most equations of state for high-density matter require a central density of $\rho_c \geq 3\rho_0$ to support a mass of $1.3-1.5\ M_\odot$, which is the range of observed neutron star masses. This is true even for equations which disregard hyperon formation (erroneously, acording to the above remarks), and is pronounced when hyperons are taken into account, since the inclusion of more baryon species softens the equation of state [16,18], calling for an even larger value of ρ_c. Hence, the most likely conclusion is that if the SQM hypothesis is correct, then all neutron stars should be strange quark stars. The conversion into strange quark matter will occur immediately at birth of the neutron star, perhaps after the initial neutrino-diffusion time. In any case, a "delayed" burning of a neutron star into a strange quark star later in its evolution seems to be ruled out.

It can be argued that the nuclear equation of state might be stiff enough so that a star of $1.4M_\odot$ will have a central density lower than $2.5\rho_0$. Such equations cannot be excluded (and are sometimes found in relativistic mean field models, due to the reduced values of effective masses in these models), although this is inconsistent with most published equations of state. However, a very stiff nuclear equation of state yields high values of the energy per baryon, and is thus susceptible to 2-flavor deconfinement. Only a very limited range of quark matter models (typically with $\alpha_c \approx 0$ and appropriate relatively high values of B) can keep the nuclear matter in stiff equations confined up to $2.5\rho_0$, while still predicting that SQM is absolutely stable. Furthermore, a stiff nuclear equation of state is also unfavorable in view of the current theory of core-collapse (type II) supernovae.

Could all neutron stars be strange quark stars? The observation of glitch behaviour in pulsars severely limits this possibility. Glitches are sudden jumps in the rotation frequency of a pulsar, with a spin down rate of $\Delta\dot{\Omega}/\dot{\Omega} \sim 10^{-3} - 10^{-2}$. Current models of the glitch phenomena rely on the neutron superfluid vortex creep theory (see [21] for a review), which requires that the effective moment of inertia of the inner crust of the star, I_i, fulfill the condition $I_i/I \approx \Delta\dot{\Omega}/\dot{\Omega}$, where I is the total moment of inertia of the star. This condition is typical of any two-component model for pulsar glitches. Following this analysis, Alpar [22] pointed out that this entire class of models for glitches must be discarded for strange quark stars, which are expected to have very small crusts ($I_i/I \sim 10^{-5}$) [3]. Glendenning and Weber have argued [23] that glitches could originate even in the very low mass crust

of strange stars, but up to date no model for strange quark star glitch has been suggested. Hence, it currently seems reasonable to conclude that at least glitching neutron stars are not strange quark stars. This argument has been used with regard to the possibility that the flux of strange quark nuggets in the universe is large enough to have converted all neutron stars to strange stars [13]. In contrast, the present Letter points to the formation of hyperons as a more robust mechanism that could convert all neutron stars to strange quark stars, and is independent of the uncertainties in estimating the rate of SQM nugget production in binary coalescence.

In conclusion, it seems likely that all neutron stars should have central densities which allow for formation of a significant hyperon fraction. This result is basically model independent and suggests, as demostrated above, a robust mechanism for the creation of strange quark matter in all neutron stars, if such matter is indeed the true ground state of baryonic matter. If this is the case, all neutron stars should convert at birth to strange quark stars - a possibility which is difficult to combine with the lack of an explanation for the observed pulsar glitch phenomena. Hence, widely accepted evaluations of hyperon formation in neutron stars [15-18] serve as an indication that strange quark matter is not the true ground state of matter.

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References

[1] E. Witten, Phys. Rev. D **30**, 272 (1984).

[2] E. Farhi and R. L. Jaffe, Phys. Rev. D **30**, 2379 (1984).

[3] C. Alcock, E. Farhi and A. Olinto, Astrophys. J. **310**, 261 (1986).

[4] P. Haensel, J. L. Zdunik and R. Schaeffer, Astron. and Astrophys. **160**, 121 (1986).

[5] G. Baym, R. Jaffe, E. W. Kolb, L. McLerran and T. P. Walker, Phys. Lett. B **160**, 181 (1985).

[6] For a review see N. K. Glendenning, *Compact Stars*, (Springer, New York, 1996).

[7] C. Schaab, B. Hermann, F. Weber, and M. K. Wiegel, Astrophys. J. Lett. **480**, L111 (1997).

[8] N. K. Glendenning, C. Kettner and F. Weber, Phys. Rev. Lett. **74**, 3519 (1995); N. K. Glendenning, C. Kettner and F. Weber, Astrophys. J. **450**, 253 (1995).

[9] A. V. Olinto, Phys. Lett. B **192**, 71 (1987).

[10] A. V. Olinto, Nucl. Phys. **B24** (Proc. Suppl.), 103 (1991).

[11] H. Heiselberg, G. Baym and C. J. Pethick, Nucl. Phys. **B24** (Proc. Suppl.), 144 (1991).

[12] K. S. Cheng and Z. G. Dai, Phys. Rev. Lett. **77**, 1210 (1996).

[13] R. R. Caldwell and J. L. Friedman, Phys. Lett. B **264**, 143 (1991).

[14] V. Thorsson, M. Prakash and J.M. Lattimer, Nucl. Phys. **A572**, 693 (1994).

[15] J. Schaffner and I.N. Mishustin, Phys. Rev. C **53**, 1416 (1996).

[16] S. Balberg and A. Gal, Nucl. Phys. A (in press, 1997), LANL preprint archive index nucl-th/9704013.

[17] N.K. Glendenning, Astrophys. J. **293** (1985) 470.

[18] P.J. Ellis, R. Knorren and M. Prakash, Phys. Lett. B **349**, 11 (1995); R. Knorren, M. Prakash and P.J. Ellis, Phys. Rev. C **52**, 3470 (1995).

[19] N. K. Glendenning, Phys. Rev. D **46**, 1274 (1992).

[20] M. Prakash, J.R. Cooke and J. M. Lattimer, Phys. Rev. D **52**, 661 (1995).

[21] D. Pines and M. A. Alpar, Nature **316**, 27 (1985).
M. A. Alpar, H. F. Chau, K. S. Cheng and D. Pines, Astrophys. J. **409**, 345 (1993).

[22] M. A. Alpar, Phys. Rev. Lett. **58**, 2152 (1987).

[23] N. K. Glendenning and F. Weber, Astrophys. J. **400**, 647 (1992).

Figure Captions

Figure. 1. Equilibrium compositions for matter containing hyperons as well as nucleons and leptons, for a model similar to the model $\delta = \gamma = \frac{5}{3}$ described in [16].

Figure. 2. Energy per baryon of baryonic matter in equilibrium, and the energy per baryon for quark matter of identical composition and density, as a function of the baryon number density. The energy per baryon in the baryonic phase (solid line) is calculated with the same model as in Figure 1, and the various curves of quark matter (dashed lines) correspond to bag models with different values of the bag constant, B, (in MeV fm^{-3}) and the coupling constant, α_c, given in brackets as (B, α_c). The arrow marks the density where the first hyperons appear ($\sim$0.3 fm^{-3}).

Figure 1

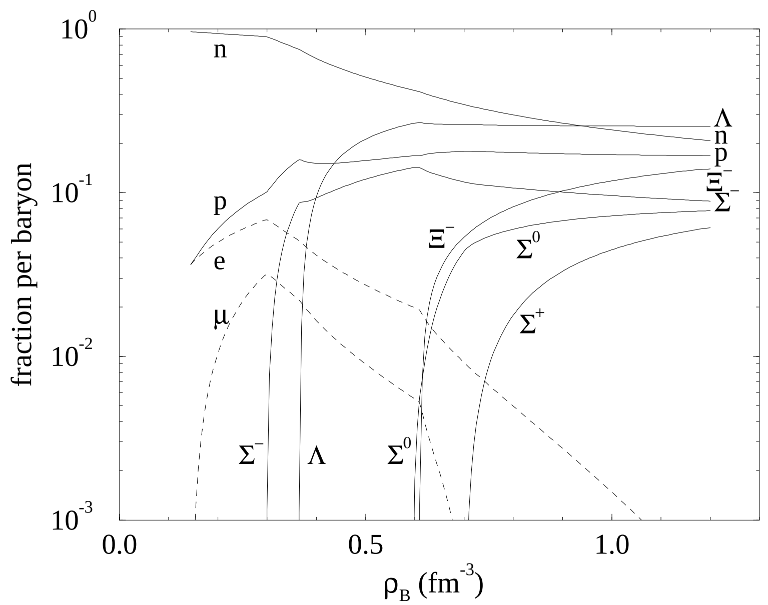

Figure 2